

Neutron Electric Dipole Moment under Non-Universal Soft SUSY Breaking Terms

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Abstract

The electric dipole moment of the neutron (EDMN) is re-examined in a general framework of the soft supersymmetry breaking. We point out some features of the relation between the EDMN and non-universal soft supersymmetry breaking terms. We give the constraints on the soft scalar masses and the soft CP phases, which have the rather large dependence on the non-universality of the soft breaking terms. We also show that the soft CP phase ϕ_B which has no natural suppression mechanism may not have large contribution to the EDMN in a certain parameter region where the radiative symmetry breaking occurs successfully. ϕ_B may not need to be so small.

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The minimal supersymmetric standard model (MSSM) is now considered as the most promising extension of the standard model (SM)[1]. Although the origin of the supersymmetry breaking is still left unknown we can make various predictions by using the suitable parametrization of its breaking. This parametrization is known as the soft supersymmetry breaking terms. Phenomenological features of the MSSM are determined by these soft supersymmetry breaking parameters, which play a similar role to a vacuum expectation value of the Higgs field in the SM. Usually these soft supersymmetry breaking parameters are treated as the universal ones as derived from a special type of supergravity theory. The electric dipole moment of the neutron (EDMN) in the MSSM is also studied under the assumption of the universal soft supersymmetry breaking terms. In such a study the EDMN is known to exceed the present experimental bound $1.1 \times 10^{-25} \text{e cm}$ [2] unless either the CP phases in the soft breaking parameters are unnaturally small of order $O(10^{-3})$ or the masses of superpartners of quarks are heavier than $O(1) \text{TeV}$ [3, 4, 5]. Both conditions seem not to be easily satisfied if we consider the soft breaking terms on the basis of the naturalness in the general way. This means that the EDMN may be a very important phenomenon to study the evidence of the supersymmetry and the origin of the supersymmetry breaking.

As recently stressed, the soft supersymmetry breaking parameters are non-universal in the effective theories which are derived from the superstring theories and also the general supergravity theories[6, 7, 8]. It is shown in some works that such a non-universality shows very interesting effects in the gauge coupling unification, the radiative symmetry breaking and so on[9, 10]. Very remarkable features which are not seen in the usual study are found in those works. The EDMN is usually very enhanced due to the presence of superpartners and new CP phases in the supersymmetric models. Therefore its re-examination under non-universal soft supersymmetry breaking terms seems to be very important for the phenomenological study of the supersymmetry and the supersymmetric model building.

In the present letter we investigate the EDMN in the MSSM with the general soft supersymmetry breaking terms and discuss its relation to the squark masses and the soft CP phases. In the MSSM it is well-known that the EDMN has non-zero value at the one-loop level due to the contributions of the gluinos, the charginos and the neutralinos. Here we concentrate ourselves only on the gluino contribution to see the effects due to the non-universality of the soft supersymmetry breaking terms to the EDMN. For

the full quantitative estimation of the EDMN we need to take account of the chargino contribution. We will comment on the chargino contribution later.

At first we briefly review the general formulae and new CP phases of the soft supersymmetry breaking terms in the MSSM and then give an explicit formula of the EDMN due to the gluino contribution. The superpotential of the MSSM is written as

$$W = h_{IJ}^U Q^I H_2 U^J + h_{IJ}^D Q^I H_1 D^J + h_{IJ}^E L^I H_1 E^J + \mu H_1 H_2, \quad (1)$$

where I and J are the generation indices. The soft supersymmetry breaking terms are

$$\begin{aligned} \mathcal{L}_{\text{soft}} = & - \sum_i \tilde{m}_i^2 |z_i|^2 + (A_{IJ}^U h_{IJ}^U Q^I H_2 U^J + A_{IJ}^D h_{IJ}^D Q^I H_1 D^J + A_{IJ}^E h_{IJ}^E L^I H_1 E^J \\ & + B\mu H_1 H_2 + \sum_a \frac{1}{2} M_a \bar{\lambda}_a \lambda_a + h.c.), \end{aligned} \quad (2)$$

where the first term represents the mass terms of all the scalar components in the MSSM. In the last term λ_a are the gaugino fields for the gauge group specified by a . The remaining terms are the scalar two and three points couplings.

Various works based on the superstring theories and also general supergravity theories suggest that these soft breaking parameters are generally non-universal[6, 7, 8]. In general low energy effective supergravity theories are characterized in terms of the Kähler potential K , the superpotential W and the gauge kinetic function f_a . Each of these is a function of the ordinary chiral matter superfields Ψ^I and the gauge singlet fields Φ^i called moduli. Usually it is assumed that the nonperturbative phenomena such as a gaugino condensation occur in the hidden sector. After integrating out the relevant fields to these phenomena, the Kähler potential and the superpotential are expanded by the low energy observable matter fields Ψ^I ,

$$K = \kappa^{-2} \hat{K}(\Phi, \bar{\Phi}) + Z(\Phi, \bar{\Phi})_{IJ} \Psi^I \bar{\Psi}^J + \left(\frac{1}{2} H(\Phi, \bar{\Phi})_{IJ} \Psi^I \Psi^J + \text{h.c.} \right) + \dots, \quad (3)$$

$$W = \hat{W}(\Phi) + \frac{1}{2} \tilde{\mu}(\Phi)_{IJ} \Psi^I \Psi^J + \tilde{h}(\Phi)_{IJK} \Psi^I \Psi^J \Psi^K + \dots, \quad (4)$$

where $\kappa^2 = 8\pi/M_{\text{pl}}^2$. The ellipses stand for terms of higher orders in Ψ^I . Using these functions the scalar potential V can be written,

$$V = \kappa^{-2} e^G [G_\alpha (G^{-1})^{\alpha\bar{\beta}} G_{\bar{\beta}} - 3\kappa^{-2}] + (\text{D-term}), \quad (5)$$

where $G = K + \kappa^{-2} \log \kappa^6 |W|^2$ and the indices α and β denote Ψ^I as well as Φ^i . The gravitino mass $m_{3/2}$ which characterizes the scale of the supersymmetry breaking is expressed

as

$$m_{3/2} = \kappa^2 e^{\hat{K}} |\hat{W}|. \quad (6)$$

In order to get the low energy effective theory from eq.(5) we take the flat limit $M_{\text{pl}} \rightarrow \infty$ with $m_{3/2}$ fixed. Through this procedure we get the superpotential (1) and the soft supersymmetry breaking terms (2). In the superpotential Yukawa couplings are rescaled as $h_{IJK} = e^{\hat{K}/2} \tilde{h}_{IJK}$ and μ term is effectively expressed as

$$\mu = e^{\hat{K}/2} \tilde{\mu} + m_{3/2} H - F^{\bar{j}} \partial_{\bar{j}} H, \quad (7)$$

where μ should be understood as $\mu_{H_1 H_2}$. From this expression we can find that the favorable μ scale of order $m_{3/2}$ can be remarkably induced. However, it should be noted that the scale of μ crucially depends on its origin and a case such as $|\mu|/m_{3/2} \ll 1$ can also occur. This case will be interesting to consider the effects of the soft CP phases on the EDMN. Each soft breaking term is expressed by using K and W as follows[7],

$$\tilde{m}_{I\bar{J}}^2 = m_{3/2}^2 Z_{I\bar{J}} - F^i \bar{F}^{\bar{j}} [\partial_i \partial_{\bar{j}} Z_{I\bar{J}} - (\partial_{\bar{j}} Z_{K\bar{J}}) Z^{K\bar{L}} (\partial_i Z_{L\bar{L}})] + \kappa^2 V_0 Z_{I\bar{J}}, \quad (8)$$

$$A_{IJ} = F^i [(\partial_i + \frac{1}{2} \hat{K}_i) h_{IJK} - Z^{\bar{M}L} \partial_i Z_{\bar{M}(I} h_{JK)L}] / h_{IJK}, \quad (9)$$

$$B = F^i [(\partial_i + \frac{1}{2} \hat{K}_i) \mu - Z^{\bar{M}L} \partial_i Z_{\bar{M}(I} \mu_{J)L}] / \mu - m_{3/2} \\ + [F^i (\partial_i + \frac{1}{2} \hat{K}_i) F^{\bar{j}} - 2m_{3/2} F^{\bar{j}}] \partial_{\bar{j}} H / \mu, \quad (10)$$

where F^i are F-terms of Φ^i and ∂_i denote $\partial/\partial\Phi^i$. V_0 in eq.(8) is the cosmological constant expressed as $V_0 = \kappa^{-2} (F^i \bar{F}^{\bar{j}} \partial_i \partial_{\bar{j}} \hat{K} - 3m_{3/2}^2)$.¹ Requiring the cosmological constant to be zero or sufficiently small, we get $|F^i| = O(m_{3/2})$. From this we find that the soft breaking terms $m_{I\bar{J}}$, A_{IJ} and B are generally non-universal but characterized by the gravitino mass $m_{3/2}$. The gaugino masses M_a are derived through the following equation,

$$M_a = \frac{1}{2} \bar{F}^{\bar{j}} \partial_{\bar{j}} \log \text{Re} f_a, \quad (11)$$

where the subscript a represents a gauge group. This shows that M_a is also characterized by $m_{3/2}$.

The values of these soft breaking terms are at most $O(m_{3/2})$. However, it should be noted that this does not mean the nonexistence of the large difference among the soft breaking parameters. In some models the soft breaking terms with the different order

¹In these formulae we do not assume that the cosmological constant vanishes.

of magnitudes are given[8, 11]. These soft breaking terms are considered as the values at M_{pl} . The non-universality at high energy scale does not necessarily mean the non-universality at the low energy when the quantum corrections are taken into account. It is also well-known that this non-universality for the scalar masses at the low energy region is strictly constrained by the FCNC phenomena[12]. In the following study we concentrate ourselves to the case where the low energy non-universality of the soft scalar masses remains without contradicting the FCNC constraints. As such a typical example, we consider the soft scalar masses in which the up and down sector squark mass matrices have the similar form and $\tilde{m}_{U_R}^2 = \tilde{m}_{D_R}^2 \neq \tilde{m}_{Q_L}^2$.² In addition we assume that $A_{IJ}^U = A_U$ and $A_{IJ}^D = A_D$ for simplicity.

The soft breaking parameters A_f , B and M_a are generally complex and then become new origins of the CP violation which do not exist in the SM. All of the phases of these parameters are known not to be physically independent. We can extract the physically independent phases from them in the usual way[4]. We make the VEVs of the Higgs fields H_1 and H_2 real by the appropriate redefinition of H_1 and H_2 so as to make B real. If we note that the complex phases in the gaugino masses are common in eq.(11), we can make the gaugino mass real by the use of the R-transformation and summarize the new CP phases associated to the soft breaking terms in the following form,

$$\phi_{A_f} = \arg(A_f M^*), \quad \phi_B = \arg(B M^*). \quad (12)$$

The non-universality of A_f terms introduces the CP phase for each Yukawa coupling. These new CP phases can cause the new contributions to the EDMN.

Now we proceed to express the formula of the EDMN explicitly keeping the non-universality of the soft breaking terms[5]. As mentioned earlier, we only consider the gluino contribution whose Feynmann diagram is shown in Fig.1. To calculate this diagram we need an explicit form of a squark mass matrix \mathcal{M}_f^2 . For the f -type squark ($f = U, D$) it is explicitly written down as

$$\begin{pmatrix} |m_f|^2 + \tilde{m}_{fL}^2 + m_Z^2 \cos 2\beta(T_f^3 - Q_f \sin^2 \theta_W) & m_f(A_f + R_f \mu^*) \\ m_f^*(A_f^* + R_f \mu) & |m_f|^2 + \tilde{m}_{fR}^2 + m_Z^2 \cos 2\beta Q_f \sin^2 \theta_W \end{pmatrix}, \quad (13)$$

² This type of soft scalar mass is discussed in ref.[9], where the interesting feature of the gauge coupling unification is pointed out.

where m_f , \tilde{m}_{fL} and \tilde{m}_{fR} are masses of the f -quark, the corresponding left-handed squark and the right-handed squark, respectively. T_f^3 is the third component of the weak isospin of left-handed quark f and Q_f is an electric charge of the quark f . R_f is defined by using $\tan \beta \equiv \langle H_2 \rangle / \langle H_1 \rangle$ as

$$R_f = \begin{cases} \cot \beta & (\text{for } T_f^3 = \frac{1}{2}), \\ \tan \beta & (\text{for } T_f^3 = -\frac{1}{2}). \end{cases} \quad (14)$$

Although \mathcal{M}_f^2 is a 6×6 matrix, here we extract the part corresponding to the first generation to estimate the EDMN. This treatment will be justified because the generation mixing off-diagonal components of \mathcal{M}_f^2 should be suppressed from the FCNC constraints.

The contribution to the EDM of a quark f from the diagram in Fig.1 is

$$d_{f/e}^g = \frac{\alpha_S}{3\pi} \sum_{i=1}^2 \text{Im}(S_{2i}S_{1i}^*) \frac{Q_f}{m_g} r_i \int_0^1 dx \frac{x(1-x)}{1-x+r_i-x(1-x)s_i} \quad (15)$$

where $r_i = m_g^2/\tilde{m}_i^2$, $s_i = m_f^2/\tilde{m}_i^2$ and m_g is the gluino mass. S_{ij} is the element of the unitary matrix S which diagonalizes the squark mass matrix \mathcal{M}_f^2 as $\mathcal{M}_{\text{diag}}^2 = S^\dagger \mathcal{M}_f^2 S$. The eigenvalues of this matrix are represented as \tilde{m}_i^2 . After evaluating these matrix elements S_{ij} , we get the final form of d_f^g as,

$$d_{f/e}^g = \frac{\alpha_S}{3\pi} \frac{Q_f \sin \gamma_f}{m_g} \left[\frac{4Z_f^2}{Y_f^2 + 4Z_f^2} \right]^{\frac{1}{2}} [r_2 I(r_2) - r_1 I(r_1)] \quad (16)$$

where \tilde{m}_1^2 and \tilde{m}_2^2 are mass eigenvalues of M_f^2 . They are explicitly written down as

$$\begin{aligned} \tilde{m}_{1f}^2 &= \frac{m_g^2}{2} \left[X_f + \sqrt{Y_f^2 + 4Z_f^2} \right], \\ \tilde{m}_{2f}^2 &= \frac{m_g^2}{2} \left[X_f - \sqrt{Y_f^2 + 4Z_f^2} \right]. \end{aligned} \quad (17)$$

Here the parameters X_f, Y_f and Z_f are defined as

$$\begin{aligned} X_f &= \frac{\tilde{m}_{fL}^2}{m_g^2} + \frac{\tilde{m}_{fR}^2}{m_g^2} + \frac{m_Z^2}{m_g^2} \cos 2\beta T_f^3, \\ Y_f &= \frac{\tilde{m}_{fL}^2}{m_g^2} - \frac{\tilde{m}_{fR}^2}{m_g^2} + \frac{m_Z^2}{m_g^2} \cos 2\beta (T_f^3 - 2Q_f \sin^2 \theta_W), \\ Z_f &= \frac{1}{m_g^2} |m_f(A_f + R_f \mu^*)|. \end{aligned} \quad (18)$$

In the derivation of these formulae we use the fact $m_f \ll m_g$ for $f = U, D$. The function $I(r)$ has the following form

$$I(r) = \frac{1}{2(1-r)^2} \left[1 + r + \frac{2r \ln r}{1-r} \right]. \quad (19)$$

In this expression $s_i = m_f^2/\tilde{m}_i^2$ is neglected because the quark mass is small enough compared to the soft breaking scalar masses. In eq.(16) an angle γ_f can be written down as³

$$\tan \gamma_f = \frac{|A_f| \sin \phi_{A_f} + |R_f \mu^*| \sin \phi_B}{|A_f| \cos \phi_{A_f} + |R_f \mu^*| \cos \phi_B}. \quad (20)$$

In the case that ϕ_{A_f} and ϕ_B are sufficiently small this reduces to the usually known form,

$$\gamma_f \sim \frac{|A_f|}{|A_f| + |R_f \mu^*|} \phi_{A_f} + \frac{|R_f \mu^*|}{|A_f| + |R_f \mu^*|} \phi_B. \quad (21)$$

To reconstruct the EDMN from ones of the quarks, we follow the conventional method and use the result of the nonrelativistic quark model

$$d_n = \frac{1}{3}(4d_d - d_u). \quad (22)$$

Now we analyze the EDMN using these formulae in the non-universal soft breaking case. In this analysis we assume $\gamma_U = \gamma_D \equiv \gamma$ and $A_U = A_D$, for simplicity. To see the effects of the non-universality in the u - and d - squark masses on the EDMN we plot the contour lines of $d_n^g/e \sin \gamma$ in the (\tilde{m}_R/m_g) - (\tilde{m}_L/m_g) plane in Figs.2 ~ 4. Each graph corresponds to the various settings of m_g and $|A_f + R_f \mu^*|$ values. Here it should be noted that in these figures the contours are drawn for $d_n^g/e \sin \gamma$ but not for the direct values of d_n^g/e . Thus for the comparison of the present results to the experimental bound we must estimate $\sin \gamma$. As is easily seen from eq.(20), $\sin \gamma$ is of order one as far as both of $\tan \phi_{A_f}$ and $\tan \phi_B$ are $O(1)$. This is independent of the values of $|A_f|$ and $|R_f \mu^*|$. In this case Figs.2 ~ 4 directly represent the value of the EDMN which can be compared to the experimental bound. If both of ϕ_{A_f} and ϕ_B are less than $O(1)$, γ will be approximately expressed by eq.(21) and $\gamma = O(\phi_{A_f}), O(\phi_B)$. Taking account of these, we can read off the conditions to satisfy the experimental bound of the EDMN from these figures. The constraints usually quoted in the universal soft breaking case seem to be rather weakened by various combined effects of the non-universal soft breaking parameters. Furthermore a suitable combination of the non-universality may present the interesting possibility for the EDMN. In particular, as seen from eq.(21), as far as $|A_f|$ and $|R_f \mu^*|$ are not the same order either ϕ_{A_f} or ϕ_B will mainly contribute to the EDMN. This may open the new possibility for the solution of the soft CP phases as seen later.

³ Here we assume that the $\arg(m_f)$ is small enough not to bring the strong CP problem.

Here we comment on the rather large dependence of the EDMN on $\tilde{m}_L, \tilde{m}_R, m_g$ and $|A_f + R_f \mu^*|$. The large values of \tilde{m}_L or \tilde{m}_R are required to suppress the EDMN. However, it should be remarkable that if the only one of them is sufficiently heavy the EDMN can be largely suppressed as expected from the consideration for the squark mass eigenvalues. As $|A_f + R_f \mu^*|$ increases $d_n^g/e \sin \gamma$ proportionally does. This feature is easily understood if we note that $|A_f + R_f \mu^*|$ characterize the left-right mixing of the squark mass matrix. The value of $d_n^g/e \sin \gamma$ increases according to the decrease of the gluino mass if we keep \tilde{m}_R/m_g and \tilde{m}_L/m_g constant. Following to the usual RGE study, the soft masses such as $m_g \gg \tilde{m}_L, \tilde{m}_R$ seem to be difficult to realize at the low energy region. The large gluino mass will make the squark masses the same order as the gluino mass through the renormalization effect.⁴ The reasonable region such as $m_g \lesssim \tilde{m}_L, \tilde{m}_R$ has the tendency to make the EDMN small. Anyway the non-universality of the soft supersymmetry breaking terms can fairly affect the value of the EDMN.

Next we study the necessity of the natural suppression of the soft CP phases. As was shown in the previous part, $\sin \gamma$ should be small enough not to exceed the experimental bound of the EDMN if all squark masses are $O(100)$ GeV. This is usually considered to be equivalent to the condition that both of the soft CP phases ϕ_{A_f} and ϕ_B are less than $10^{-2} \sim 10^{-3}$ depending on m_g and $|A_f + R_f \mu^*|$. From the viewpoint of the naturalness such small phases seem not to be expected in the general soft supersymmetry breaking schemes. In the following parts we shall propose a natural explanation for this problem.

Recently it is suggested that the phase ϕ_{A_f} can be small enough not to contradict with the EDMN bound in the models derived from the superstring theories with the supersymmetry breaking due to the F-terms of a dilaton and moduli. In ref.[8] it was shown that the dilaton dominated supersymmetry breaking suppresses the phase ϕ_{A_f} sufficiently. This is because the phase structure of A_f and M_a in eqs.(9) and (11) has a certain similarity. On the other hand Choi pointed out in ref.[13] that the various complex phases contributing ϕ_{A_f} are tuned to the value less than $O(10^{-3})$ by the dynamical mechanism based on the Peccei-Quinn symmetry on the dilaton and moduli. Unfortunately there is no such general suppression mechanisms for the phase ϕ_B . The origin of μ term is

⁴ In the present study we consider only the gluino contribution and thus the non-universality of the gaugino masses is irrelevant. However, when we estimate the chargino contribution it should be also taken into account.

not determined as shown in eq.(7) and the structure of ϕ_B completely depends on its origin as seen from eq.(10). It seems very difficult to suppress ϕ_B naturally.⁵ However, the existence of natural suppression mechanisms of ϕ_{A_f} can open the new possibility of the sufficient suppression of the EDMN. Instead of finding the suppression mechanism of ϕ_B , it seems more promising to investigate this new possibility that the EDMN may be sufficiently suppressed even if the phase ϕ_B is not so small.

For this purpose we will consider the case of $|A_f| \gg |R_f\mu|$. We assume that the smallness of ϕ_{A_f} is guaranteed by the above mentioned mechanism. In such a case the value of γ_f can be estimated as

$$\gamma_f \sim \phi_{A_f} + \frac{|R_f\mu^*|}{|A_f|} \sin \phi_B. \quad (23)$$

The contribution from ϕ_B can be suppressed by a factor $|R_f\mu^*|/|A_f|$ even if ϕ_B is $O(1)$. The main issue of this scenario is the consistency between the radiative symmetry breaking and the smallness of $|R_f\mu^*|/|A_f|$. Using eqs.(16) and (23), we can estimate the contribution to the EDMN from ϕ_B as

$$d_n^g/e \sin \phi_B = \frac{1}{3} \frac{|\mu|}{|A|} (-X_U \cot \beta + 4X_D \tan \beta) \sim \frac{4}{3} \frac{|\mu|}{|A|} X_D \tan \beta \quad (24)$$

where A_f is assumed as $A_U = A_D = A$. The approximate value of $4X_D/3$ can be read off from Figs.2 ~ 4. Generally $\tan \beta > 1$ is expected in the radiative symmetry breaking scenario due to the large top Yukawa coupling. Thus the second similarity in eq.(24) is deduced. As far as $\tan \beta \sim O(1)$ and the masses of all superpartners are ~ 100 GeV, the necessary condition to satisfy the experimental bound of the EDMN is estimated as $|\mu|/|A| < 10^{-2} \sim 10^{-3}$. As suggested by the previous argument on the soft breaking terms, A is expected as $O(m_{3/2})$ where the magnitude of $m_{3/2}$ is dependent on the supersymmetry breaking mechanism. On the other hand the scale of μ depends on its origin and then $|\mu|/|A| < 10^{-2} \sim 10^{-3}$ may be possible. However, the small μ may yield the light chargino and conflict the experimental constraint. The chargino mass matrix has the following form

$$\begin{pmatrix} \mu & \sqrt{2}m_W \sin \beta \\ \sqrt{2}m_W \cos \beta & M_2 \end{pmatrix} \quad (25)$$

⁵ An only known mechanism to suppress ϕ_B naturally is to replace the μ -term by a Yukawa coupling of a singlet field and Higgs fields. In this model ϕ_B is reduced to the ϕ_A type phase and then ϕ_B will also be small automatically[8].

and then the squared mass eigenvalues of the charginos are

$$m_{\omega_{1,2}}^2 = \frac{1}{2} \left[\mu^2 + M_2^2 + 2m_W^2 \sin 2\beta \mp \sqrt{(\mu^2 - M_2^2)^2 + 4m_W^2(\mu + M_2)^2 \sin 2\beta} \right]. \quad (26)$$

If we consider the region where both of μ and M_2 are sufficiently small and also $\tan \beta$ is not so larger than one, the chargino mass does not conflict the present experimental bound 45 GeV.

Based on these considerations we concentrate ourselves to the parameter region such as

$$B, A, \tilde{m}_i \gg \mu, M_2 \quad (27)$$

at the low energy region. In this region we practice the RGE study to examine the radiative $SU(2) \times U(1)$ breaking and estimate the top quark mass. As is well known, the small μ increases the value of $\tan \beta$ significantly. This effect may cancel the smallness of μ and make our scenario less attractive. To avoid such situation, we take $B(M_{\text{pl}})$ larger than other soft parameters. This is because $\tan \beta$ does not depend on μ directly but depends on $B\mu$. At the tree level analysis, β is expressed as[14]

$$\sin 2\beta = \frac{2B\mu}{m_{H_1}^2 + m_{H_2}^2 + 2\mu^2}. \quad (28)$$

From this one can find that the value of m_{H_1} also influences $\tan \beta$ in the same way as B ⁶. The non-universality of soft scalar masses may be applied to not only squarks and sleptons but also the Higgs sector.⁷ By choosing $m_{H_1}(M_{\text{pl}})$ smaller than $m_{H_2}(M_{\text{pl}})$, we can reduce $\tan \beta$ furthermore. The non-universality between \tilde{m}_U and \tilde{m}_D also affects the running of Higgs masses through RGEs and one can expect similar effects as above. However, such effects are indirect and negligible unless we assume extremely large squark mass hierarchy, which often causes color $SU(3)$ breaking. Combining these effects we can find a suitable parameter region on the basis of RGEs study. As such an example we list a typical set of soft supersymmetry breaking parameters at m_Z ,

$$|\mu|/|A| = 2.9 \times 10^{-2}, \quad \tan \beta = 2.6.$$

⁶In the most case of the radiative symmetry breaking, $|m_{H_2}^2 + \mu^2|$ is about $O(10^{-1})(m_{H_1}^2 + \mu^2)$ so that we need not to take account of its effect.

⁷ From the viewpoint of radiative $SU(2) \times U(1)$ breaking, it is interesting to vary initial values of Higgs masses from other scalar ones[10].

For these values the top quark mass becomes 148 GeV and the lighter chargino mass is ~ 45 GeV. The constraint on ϕ_B seems to be largely reduced to the rather natural value as $\phi_B \sim O(10^{-1})$ for these parameters. The combined effects of the non-universality of the soft supersymmetry breaking parameters are expected to weaken the constraints from the EDMN furthermore.

Finally we should comment on the chargino contribution. In this case the ϕ_A dependence is largely suppressed due to the small Yukawa couplings. The chargino contribution mainly comes from the ϕ_B dependent part of the d -quark EDM and is expressed as[5]

$$\begin{aligned} d_d^c/e \sin \phi_B &\lesssim \frac{\alpha_{em}}{4\pi \sin^2 \theta_W} \frac{M_2 |\mu|}{m_{\omega_2}^2 - m_{\omega_1}^2} \frac{m_f}{\tilde{m}_i^2} \times O(1) \\ &\sim 5.1 \times 10^{-25} \left(\frac{1 \text{ TeV}}{\tilde{m}_i} \right)^2 \left(\frac{m_f}{10 \text{ MeV}} \right) \frac{M_2 |\mu|}{m_{\omega_2}^2 - m_{\omega_1}^2} \times O(1), \end{aligned} \quad (29)$$

In the parameter region presented above as an example this chargino contribution is expected to be sufficiently within the experimental bound for $\phi_B \sim 10^{-1}$ even if $m_i \sim O(100)$ GeV. This is because the appearance of the additional suppression factor

$$\frac{M_2 |\mu|}{m_{\omega_2}^2 - m_{\omega_1}^2} \sim \frac{|\mu|}{2m_W \sin^{1/2} 2\beta} \sim 7.5 \times 10^{-2}$$

for the above parameters. The present experimental bound of the EDMN may be reconciled with the soft supersymmetry breaking without introducing unnatural assumptions on the parameters. Although the present parameter region does not seem so wide, it may inform us the suitable non-universality of the soft breaking parameters in the MSSM.

In summary, we re-examined the EDMN under the general soft supersymmetry breaking parameters. We pointed out some features of the relations between the EDMN and soft supersymmetry breaking parameters which seems not to be mentioned explicitly before in the universal soft breaking framework. We also showed that the soft CP phase ϕ_B whose natural suppression mechanism is not known up to now does not have large contribution to the EDMN in the certain parameter space where the radiative symmetry breaking occurs successfully. This may be an interesting non-universal parameter region of the MSSM. We may not need to require that ϕ_B is so small. FCNC constrains severely the soft masses of the squarks, in particular, with the same charge. It requires their degeneracy at m_Z scale. On the other hand the study of the EDMN may give us some other knowledge for the squark masses as is shown in this paper. If we combine these, we may get a useful insight for the whole structure of the soft squark masses. Moreover,

the EDMN may give us useful information of the soft breaking parameters in the MSSM. From this point of view the more precise theoretical study of the EDMN seems to be very important. Also the improvement of the experimental bound of the EDMN is strongly desired.

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Figure Captions

Fig.1

A Feynmann diagram of the gluino contribution to the EDMN.

Fig.2

The contours of $d_n^g/e \sin \gamma$ in the (\tilde{m}_R/m_g) - (\tilde{m}_L/m_g) plane at $m_g = 100$ GeV. $|A_f + R_f \mu^*|$ is chosen as 100 GeV and 1000 GeV in Fig.2A and Fig.2B, respectively. Each contour corresponds to a) 2.0×10^{-26} cm, b) 2.0×10^{-25} cm, c) 2.0×10^{-24} cm, d) 2.0×10^{-23} cm and e) 2.0×10^{-22} cm.

Fig.3

The contours of $d_n^g/e \sin \gamma$ in the (\tilde{m}_R/m_g) - (\tilde{m}_L/m_g) plane at $m_g = 500$ GeV. The setting of $|A_f + R_f \mu^*|$ is the same as Fig.2. Each contour represents the same value as Fig.2.

Fig.4

The contours of $d_n^g/e \sin \gamma$ in the (\tilde{m}_R/m_g) - (\tilde{m}_L/m_g) plane at $m_g = 1000$ GeV. The setting of $|A_f + R_f \mu^*|$ is the same as Fig.2. Each contour represents the same value as Fig.2. The region outside the vertical and horizontal lines is prohibited because the squared mass eigenvalues of squark become negative.

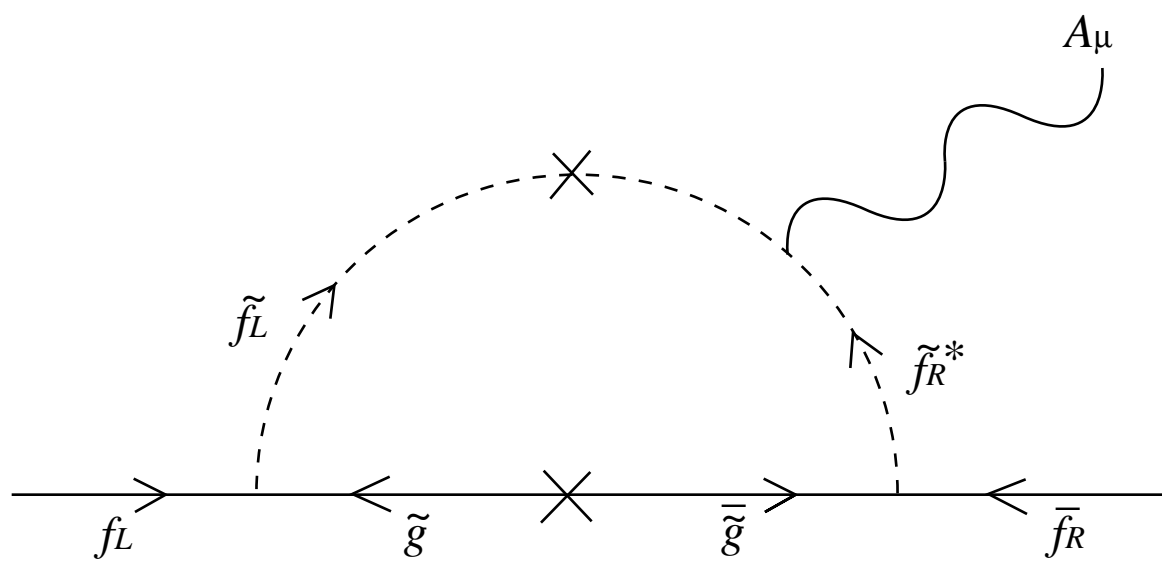


Fig.1

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Fig.2-A

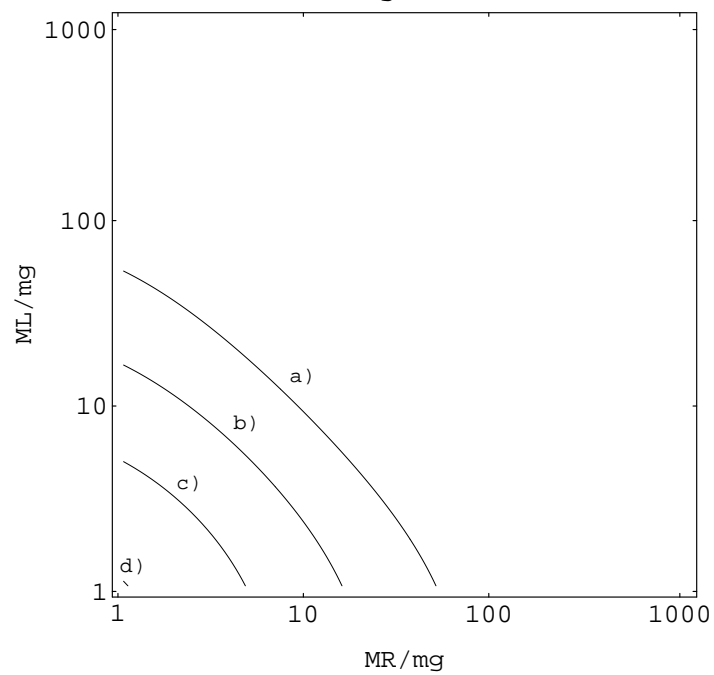
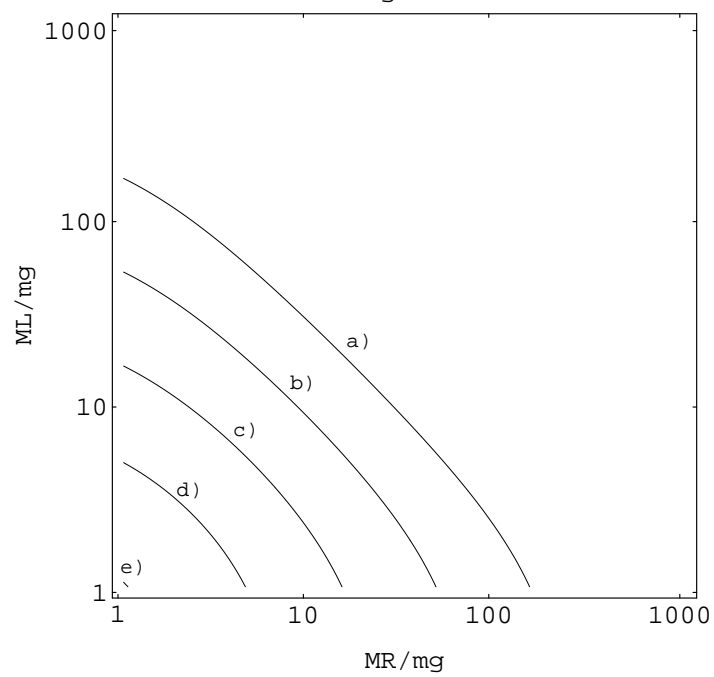


Fig.2-B



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Fig.3-A

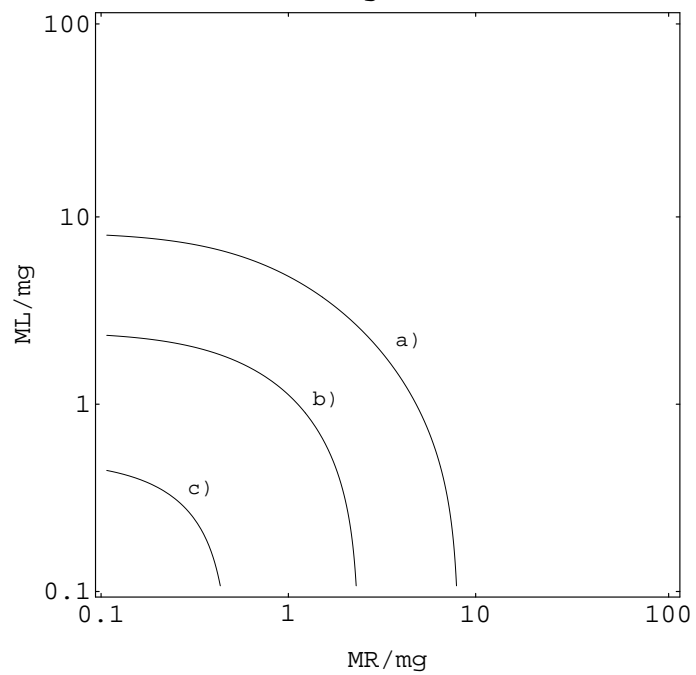
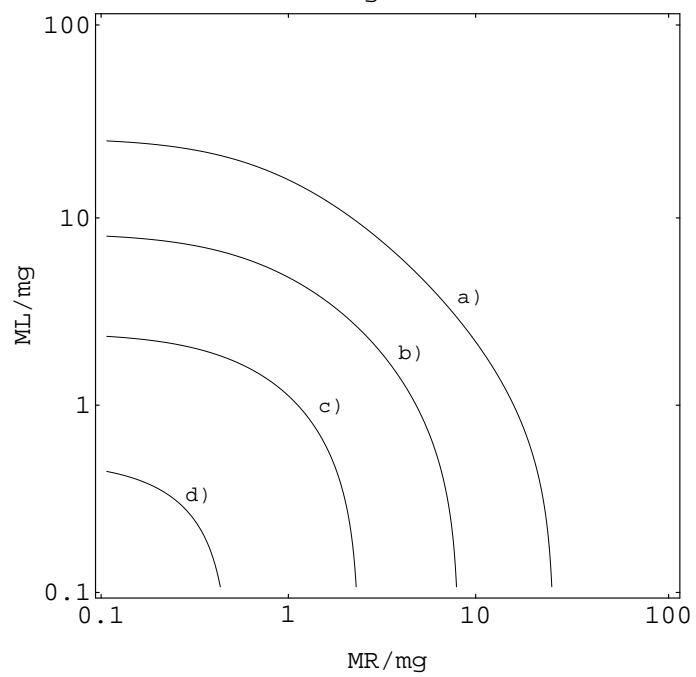


Fig.3-B



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Fig.4-A

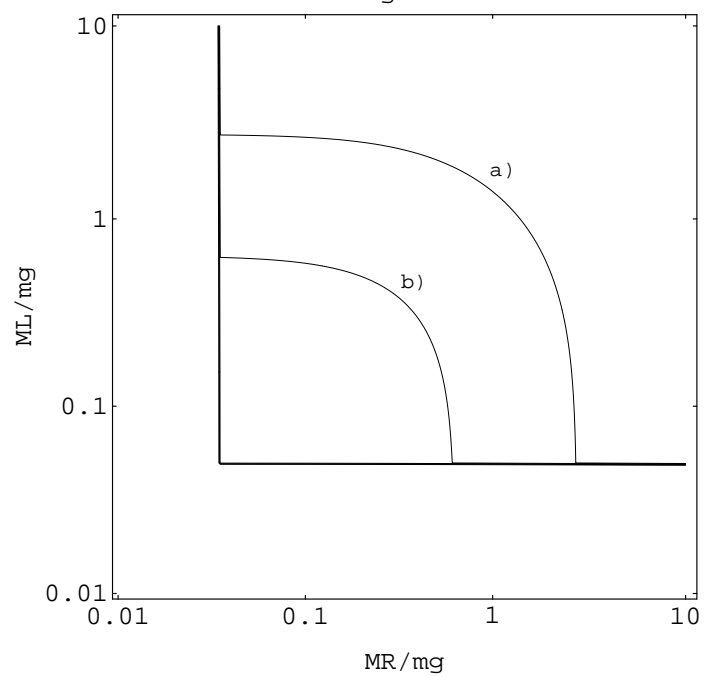
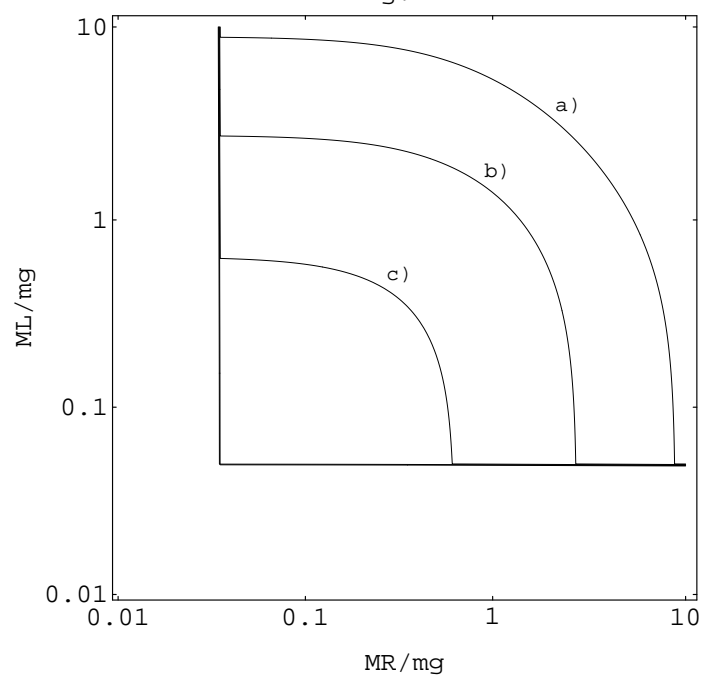


Fig.4-B



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